“Gaugomaly” Mediated SUSY Breaking and Conformal Sequestering

Raman Sundrum*

Department of Physics and Astronomy
Johns Hopkins University
3400 North Charles St.
Baltimore, MD 21218-2686

Abstract

Anomaly-mediated supersymmetry breaking in the context of 4D conformally sequestered models is combined with Poppitz-Trivedi D-type gauge-mediation. The implementation of the two mediation mechanisms naturally leads to visible soft masses at the same scale so that they can cooperatively solve the μ and flavor problems of weak scale supersymmetry, as well as the tachyonic slepton problem of pure anomaly-mediation. The tools are developed in a modular fashion for more readily fitting into the general program of optimizing supersymmetric dynamics in hunting for the most attractive weak scale phenomenologies combined with Planck-scale plausibility.

*email: sundrum@pha.jhu.edu

1 Introduction

The scenario of weak scale supersymmetry (SUSY) (reviewed in Ref. [1]), while broadly attractive, faces several technical challenges: A successful theory should explain (i) what mechanism robustly protects rare processes such as flavor-changing neutral currents (FCNCs) from excessive superpartner-mediated corrections, (ii) how gauge coupling unification, apparent in the softly broken minimal supersymmetric standard model (MSSM), is preserved in the full theory, (iii) how visible sector SUSY breaking triggers electroweak symmetry breaking, and (iv) how all superpartners and Higgs bosons naturally obtain masses that evade present search constraints. While economical mechanisms are available to address any of these issues individually, addressing them collectively is rather difficult. The purpose of this paper is to extend the set of mechanisms upon which realistic model-building may be based. Indeed a realistic model is presented, although it is hoped that further investigation, and perhaps additional mechanisms, will yield models with larger acceptable parameter spaces.

Issue (i), in particular the SUSY flavor problem, is sensitive to the mechanism by which SUSY breaking, originating in a hidden sector, is mediated to the visible sector. A promising approach is to have the mediating force be flavor-blind, the natural candidates being gauge forces or gravity. Gauge-mediated SUSY breaking (GMSB) [2] [3] (reviewed in Ref. [4]), as well as gaugino-mediated SUSY breaking [5], are based on the first possibility, while anomaly-mediated SUSY breaking [6] [7] is based on the second. While GMSB and AMSB are compatible with (ii), they both face significant difficulties with (iii) and (iv). Both suffer from the μ-problem. AMSB applied to the MSSM yields tachyonic sleptons [6] which destabilize the standard electroweak vacuum. In this paper, we will discuss how GMSB and AMSB can naturally combine symbiotically in addressing (i) – (iv). This general direction is not new. Ref. [8] demonstrated within a non-minimal visible sector that AMSB effects could trigger competitive GMSB contributions. Ref. [9] proposed "gaugino-assisted" AMSB. Refs. [10] [11] combined AMSB with SUSY breaking arising from gauge D-terms. Finally, Refs. [12] [13], which are of particular relevance to the present paper, proposed that AMSB could naturally combine with the "D-type" of GMSB studied by Poppitz and Trivedi [14].

A central issue is that of sequestering, ensuring that the AMSB and GMSB contributions to the visible sector are not overshadowed by SUSY breaking from Planck-suppressed (superspace) contact interactions directly with the hidden sector. Originally this was achieved by considering the visible and hidden sectors to be separated in a minimally inhabited extra dimension [6] [15]. While this setting is quite economical and technically natural, it does introduce the extra element of non-renormalizable extra-dimensional effective field theory on top of non-renormalizable 4D (super-)gravity. While string theory strongly motivates the existence of extra

dimensions, it is not easy to assess how readily extra dimensions with the requisite minimal light field content arise. Refs. [16] studied string theoretic examples where these conditions were violated, but it is still hard to know how generic their conclusions are.

A more recent approach to sequestering is to remain in 4D, but dynamically suppress visible-hidden contact interactions through strong renormalization effects in running them down from the Planck scale. This requires either the visible sector or the hidden sector to be strongly coupled over a large hierarchy (and therefore approximately conformal) below the Planck scale. Refs. [17] studied the first possibility in the context of gaugino-mediated SUSY breaking. Refs. [18] [19] studied the latter possibility as well as the conceptual connections between this conformal sequestering and extra-dimensional sequestering (in particular of the warped form studied in Ref. [20]) via the AdS/CFT correspondence [21] [22] [23]. The advantage here is that, when based on known super-conformal field theories (SCFT's), there is no doubt as to the *existence and properties* of the sequestering dynamics.

Conformal sequestering has been criticized recently [24] for its reliance on hidden discrete symmetries, or non-genericity in order to avoid reliance on these symmetries. It is still unclear, because of our relative ignorance of the whole "landscape" of strongly coupled SCFT's, whether such aesthetic considerations reflect a serious and generic problem with conformal sequestering or our restriction thusfar to simple controllable examples.

The present paper will employ the conformal sequestering dynamics of Ref. [19], used there in a pure AMSB context, but extended here with a messenger sector that leads to competitive GMSB contributions to the visible sector[1]. The results particularly resemble those of the extra-dimensionally sequestered models of Ref. [12], although the mechanisms are quite distinct (even after mapping qualitative features via the AdS/CFT correspondence). It should also be mentioned that realistic models employing conformal sequestering, combining AMSB and D-terms of gauged $B-L$ and hypercharge were developed in Ref. [25].

To appreciate the level of theoretical control, note that the effective interactions presented below do employ non-renormalizable (but 4D) Planck-suppressed interactions, which are natural to expect along with gravity. But, with the obvious exception of (super-)gravity, all of the interactions can be made renormalizable at the expense of introducing some extra massive fields. We will not do this here. The important physical scales will lie far below M_{Pl}, so that only classical supergravity is important. That is, the level of control is similar to doing QED in the earth's gravitational field.

[1]The particular hidden sector model in Ref. [18] depends on a certain strong interaction coefficient being negative, whereas its sign can in fact be determined to be positive by chiral lagrangian symmetry arguments. The model can be remedied, and this will be detailed elsewhere. We proceed here with Ref. [19] which avoids this problem.

The organization of this paper is as follows. The salient features of Ref. [19] will be reviewed in Section 2. A compatible messenger sector will be added in Section 3, explaining how GMSB is naturally competitive with AMSB. Section 4 will present a realistic model (except that we have not bothered with neutrino masses) based on having the visible sector be the next-to-minimal SUSY standard model (NMSSM). Section 5 discusses extensions aimed at broadening the parameter space in different directions, chiefly to give maximal freedom to the energy scales for the *origins* of visible flavor structure, while still solving the supersymmetric flavor problem.

2 Hidden Sector Dynamics

We will adopt the hidden sector construction of Ref. [19], which we therefore briefly review in this section.

2.1 Dynamical SUSY breaking

The hidden sector comprises two SQCD sub-sectors with Planck-suppressed couplings between them. One of the sub-sectors, $SQCD_2$, has two colors and four "flavors" (eight doublets), denoted for convenience by $T_{Ia}, I = 1, 2, 3, 4, a = 1, 2$. The theory lies at the self-dual point of Seiberg's conformal window, and flows to a strongly-coupled IR fixed point. We assume that at the Planck scale this subsector is already near this fixed point. The other sub-sector, $SQCD_3$, has three colors and two flavors, $P_a, \overline{P}_a,\ a = 1, 2$, and is weakly coupled at the Planck scale. The hidden superpotential is given by

$$W_{hid} = \frac{\lambda'}{M_{Pl}} \sum_{I \neq J} (T_{aI} T_I^a)(T_{bJ} T_J^b) + \frac{\lambda}{M_{Pl}} \sum_I (T_{aI} T_I^a)(\overline{P}_a P^a), \tag{2.1}$$

where $\lambda, \lambda' \sim \mathcal{O}(1)$, the parantheses denote gauge-invariant bilinears, and the a index is raised using the ϵ_{ab}-tensor.

There are regions of the *classical* hidden moduli space where it is one(-complex)-dimensional, parametrized by an $SQCD_2$ quark bilinear, $(T_{1a} T_1^a)$. However, non-perturbative effects lift and stabilize this modulus. First, because of the non-trivial fixed point dynamics of the $SQCD_2$ subsector, the canonically normalized modulus is given by $X \equiv (T_{1a} T^{1a})^{2/3}/M_{Pl}^{1/3}$, that is the effective Kahler potential at the fixed point is $\sim X^\dagger X$ for $X \ll M_{Pl}$. Secondly, non-perturbative effects from the $SQCD_3$ sub-sector, acting through λ dynamically generate a linear superpotential for X. The net result is an effective Polonyi model of SUSY breaking [26] by X, where $\langle X \rangle \ll M_{Pl}$ (unlike the original Polonyi model) is stabilized by a combination of SUGRA corrections and the leading Kahler corrections for X (determined by the flow to the

non-trivial fixed point). It should be noted that the vacuum so obtained is only metastable, but can easily be stable on cosmological time scales.

The hidden SUSY breaking is given by

$$F_X \equiv \Lambda_{int}^2, \tag{2.2}$$

where the intermediate scale $\Lambda_{int} \ll M_{Pl}$ naturally due to its non-perturbative origins. Once the cosmological constant is cancelled by inclusion of a constant superpotential, one finds

$$F_{SUGRA} \sim m_{3/2} \sim \frac{F_X}{M_{Pl}} = \frac{\Lambda_{int}^2}{M_{Pl}}. \tag{2.3}$$

2.2 Conformal sequestering, hidden symmetries and AMSB

Visible sector SUSY breaking can now arise by couplings to F_{SUGRA} and F_X. The SUGRA couplings of the visible sector ensure that SUSY breaking inherited from F_{SUGRA} is proportional to the breaking of visible conformal invariance. For a classically scale-invariant visible sector with the standard loop-level scale anomaly, the visible SUSY breaking pattern is given by AMSB [6] [7]. To compare later with gauge-mediated soft terms let us note that the visible AMSB soft masses are all of the rough form,

$$m_{soft} \sim (\alpha_{vis}/4\pi) F_{SUGRA}, \tag{2.4}$$

a visible loop factor being necessary in order to be sensitive to the quantum scale anomaly.

Visible SUSY breaking from F_X goes through effective superspace operators which mix the hidden sector with the visible sector. Such effects can dominate over AMSB and ruin the solution it offers to the SUSY flavor problem. Some of these effective visible-hidden operators (in particular superpotential terms) can naturally be avoided by appealing to non-renormalization theorems. The most dangerous unprotected operators are Kahler terms of the form

$$\Delta K \sim \frac{Q_i^\dagger Q_j T_{Ia}^\dagger T_{Jb}}{M_{Pl}^2}, \tag{2.5}$$

where i, j label visible flavors and I, J, a, b label hidden flavors. In the IR, $T_{Ia}^\dagger T_{Jb}$ becomes proportional to $X^\dagger X$ and results in the direct communication of hidden SUSY breaking to the visible sector. Since the couplings have no reason to preserve visible flavor symmetry (already broken in the standard model), visible flavor will be violated by SUSY breaking originating from F_X.

The central subtlety in estimating how $T_{Ia}^\dagger T_{Jb}$ reduces to $X^\dagger X$ in the IR is taking account of the strong renormalization effects of the $SQCD_2$ dynamics as the operator is run down from

the Planck scale down to $\langle X \rangle \ll M_{Pl}$ (which spontaneously breaks the approximate conformal invariance of the $SQCD_2$ sub-sector). Let us decompose the hidden bilinear into a hidden-flavor singlet part and a non-singlet part (containing no singlet),

$$T^\dagger_{Ia} T_{Jb} = (T^\dagger_{Ia} T_{Jb})_{non-singlet} + (T^\dagger_{Ia} T_{Jb})_{singlet}. \tag{2.6}$$

Now the non-singlet piece is a non-anomalous hidden flavor current (super-multiplet) and therefore is not renormalized by the $SQCD_2$ dynamics. Since $X^\dagger X$ is not pure singlet, couplings of the form

$$\Delta K \sim \frac{Q^\dagger_i Q_j \; (T^\dagger_{Ia} T_{Jb})_{non-singlet}}{M^2_{Pl}}, \tag{2.7}$$

translate in the IR to

$$\Delta K_{eff} \sim \frac{Q^\dagger_i Q_j X^\dagger X}{M^2_{Pl}}, \tag{2.8}$$

implying flavor-violating SUSY breaking for visible scalars,

$$m^2_{scalar} \sim \frac{|F_X|^2}{M^2_{Pl}} = \frac{\Lambda^4_{int}}{M^2_{Pl}}. \tag{2.9}$$

Such large flavor-violating SUSY breaking would make AMSB irrelevant and introduce the SUSY flavor problem. Fortunately, the visible coupling to the hidden non-singlet can be forbidden by insisting on some amount of hidden flavor symmetry. It is easy to check that W_{hid} is compatible with having exact dynamical symmetries comprising the discrete symmetries,

$$\begin{aligned} (i)\; T_{1a} &\leftrightarrow T_{2a},\; T_{3a} \leftrightarrow T_{4a} \\ (ii)\; T_{1a} &\leftrightarrow T_{3a},\; T_{2a} \leftrightarrow T_{4a} \\ (iii)\; T_{Ia} &\rightarrow -T_{Ia},\; T_{Jb} \rightarrow T_{Jb},\; J \neq I, \end{aligned} \tag{2.10}$$

as well as an $SU(2)$ symmetry acting on all a-indices[2]. This symmetry is powerful enough to accidentally forbid the visible couplings to the non-singlet hidden bilinear.

On the other hand, there is no exact dynamical symmetry that can forbid the visible sector couplings to the hidden flavor singlet bilinear,

$$\Delta K \sim \frac{Q^\dagger_i Q_j |T|^2}{M^2_{Pl}}, \tag{2.11}$$

simply because $|T|^2$ is of the form of the T kinetic term itself. The whole point of conformal sequestering is that $|T|^2$ corresponds to an anomalous current, and has an order one anomalous dimension suppressing the mixed coupling in the running between M_{Pl} and $\langle X \rangle$, quite plausibly enough to render AMSB so dominant as to solve the SUSY flavor problem.

[2]This $SU(2)$ symmetry can either be a global symmetry, or to protect it from gravitational corrections, it can be very weakly gauged. Alternatively, one only needs a non-abelian discrete subgroup of the $SU(2)$.

3 Messenger Sector

We now add a messenger sector to facilitate GMSB.

3.1 Fields, superpotential and hidden symmetries

The messenger sector consists of new chiral fields charged under standard model gauge symmetry, $\psi, \chi, \overline{\psi}, \overline{\chi}$ and singlets A, B, C, D. In order to maintain the unification of couplings of the SUSY standard model, under the embedding of the standard model gauge group into $SU(5)$, we take ψ, χ to be 5's and $\overline{\psi}, \overline{\chi}$ to be $\overline{5}$'s. The messenger sector is distinguished from the visible sector in that, while the visible sector does not transform under the hidden discrete symmetries, Eq. (2.10), the messenger fields transform non-trivially under (i),

$$(i)\ \chi \leftrightarrow \psi, \quad \overline{\chi} \leftrightarrow \overline{\psi}, \quad A \leftrightarrow B, \quad C \to C, \quad D \to D, \tag{3.1}$$

but are inert under (ii), (iii).

The messenger superpotential (compatible with the hidden discrete symmetries) is given by

$$\begin{aligned} W_{mess.} &= \lambda_+(\overline{\chi}\chi + \overline{\psi}\psi)(A+B) \ + \ \lambda_-(\overline{\chi}\chi - \overline{\psi}\psi)(A-B) \\ &+ \ \lambda_C C(A^2 + B^2 - v_{mess}^2) \ + \ \lambda_D DAB, \end{aligned} \tag{3.2}$$

with the dimensionless λ couplings again being order one. The mass parameter $v_{mess} \ll M_{Pl}$ is treated here as an input scale although it is simple to realize the small ratio v_{mess}/M_{Pl} as arising from naturally small non-perturbative effects[3].

F-flatness in the purely singlet part of the superpotential results in the VEVs

$$A = v_{mess}, \quad B = C = D = 0, \tag{3.3}$$

with all masses $\sim v_{mess}$. That is, the hidden discrete symmetry (i) is broken in the IR by v_{mess}. This results in messenger mass terms of order v_{mess},

$$W_{mess.-mass} = (\lambda_+ + \lambda_-)v_{mess}\ \overline{\chi}\chi + (\lambda_+ - \lambda_-)v_{mess}\ \overline{\psi}\psi, \tag{3.4}$$

which do not respect the hidden discrete symmetry (i). We will see why this is important in the next section.

It seems possible that the messenger and hidden sectors can be more economically combined, in particular since the messenger singlets are, with some contrivance, spontaneously breaking hidden discrete symmetries which are already broken by the hidden sector. However, in the interests of modularity this will not be pursued here.

[3]For example, a pure SUSY Yang-Mills sector coupling to C according to $\int d^2\theta\ C\mathcal{W}_\alpha^2/M_{Pl}$ will generate a non-perturbative linear superpotential (and negligible higher order terms) for C upon gaugino-condensation.

3.2 Poppitz-Trivedi Gauge Mediation

Our analysis now will be similar to some of the steps taken in Ref. [12]. Just as with the visible sector, there are Kahler couplings of the messengers to the hidden flavor singlet current, such as

$$\begin{aligned}\Delta K \quad &\sim \quad \frac{(\chi^\dagger\chi + \psi^\dagger\psi)(|T_{1a}|^2 + |T_{2a}|^2 + |T_{3a}|^2 + |T_{4a}|^2)}{M_{Pl}^2}, \\ &\quad\;\; \frac{|C|^2(|T_{1a}|^2 + |T_{2a}|^2 + |T_{3a}|^2 + |T_{4a}|^2)}{M_{Pl}^2},\end{aligned} \tag{3.5}$$

but just as in the visible case, these are conformally sequestered and insignificant. However, the non-trivial transformation of the messengers under (i) also allows one to write Kahler couplings to hidden flavor non-singlet currents, such as

$$\Delta K \quad \sim \quad \frac{(\chi^\dagger\chi - \psi^\dagger\psi)(|T_{1a}|^2 - |T_{2a}|^2 + |T_{3a}|^2 - |T_{4a}|^2)}{M_{Pl}^2}, \tag{3.6}$$

which is not renormalized by the strong $SQCD_2$ dynamics, and matches to an IR coupling of the form

$$\Delta K_{eff} \sim \frac{(\chi^\dagger\chi - \psi^\dagger\psi)X^\dagger X}{M_{Pl}^2}. \tag{3.7}$$

This yields the following soft mass-squareds for messenger scalars upon picking out the $F_X^\dagger F_X$ term[4]:

$$m^2_{\overline{\chi}soft} = -m^2_{\overline{\psi}soft} \sim m^2_{\chi soft} = -m^2_{\psi soft} \sim \frac{\Lambda_{int}^4}{M_{Pl}^2}. \tag{3.8}$$

It is this soft SUSY breaking carried by the messengers which is mediated to the visible sector scalars by their shared gauge interactions at two loops, in the manner determined by Poppitz and Trivedi [14]:

$$\begin{aligned}m^2_{vis.\ soft} \quad &\sim \quad (\alpha_{gauge}/4\pi)^2 \;\ln|\frac{\lambda_+ + \lambda_-}{\lambda_+ - \lambda_-}|\; m^2_{\chi soft} \\ &\sim \quad (\alpha_{gauge}/4\pi)^2 \;\ln|\frac{\lambda_+ + \lambda_-}{\lambda_+ - \lambda_-}|\; \frac{\Lambda_{int}^4}{M_{Pl}^2}.\end{aligned} \tag{3.9}$$

Note several significant features of the gauge-mediated contributions to the visible soft terms. *The gauge mediated visible soft masses are automatically parametrically of the same*

[4]It is straightforward to check that all terms involving the lowest component VEV $\langle X\rangle \ll M_{Pl}$ lead to subdominant effects. Of course there are messenger soft terms arising from F_{SUGRA} as well, but the resulting contributions to visible soft masses is additive to gauge-mediation, and because of the UV insensitivity of the former, is subsumed in the AMSB effects at the weak scale. These, having been reviewed in the previous section, are neglected in this section.

size as AMSB contributions, namely visible-loop-factor$\times\Lambda_{int}^2/M_{Pl}$. The gauge-mediated contributions are to visible scalar mass-squareds only, not to fermion masses, A-terms or B-terms. The sign and the precise magnitude of the contributions is a free parameter from the effective field theory point of view, being determined by the independent coefficient of the hidden non-singlet term in Eq. (3.6) (which explains why only estimates have been given for the soft terms). Thus gauge-mediated contributions can naturally solve the tachyonic-slepton problem of minimal AMSB by choice of this coefficient [12]. The 2-loop relation Eq. (3.9) is UV finite, potential log divergences cancelling because of the suppressed messenger supertrace, suppressed because of conformal sequestering. The log in Eq. (3.9) is therefore automatically order one, so that no large logs upset the balance of gauge-mediated and AMSB visible soft terms. The whole purpose of the messenger sector singlets was to spontaneously break the hidden discrete symmetries in the IR, allowing $\lambda_- \neq 0$, so that the log does not vanish.

The messenger supersymmetric mass scale, v_{mess} does not appear explicitly in Eq. (3.9) (having cancelled out of the ratio of messenger masses in the log), so it does not have to be tuned with respect to the weak scale, in principle it can lie within a very large range. We will put an upper bound on it below so that it does not spoil sequestering in the visible sector. Another consideration is that gauge mediation only solves the SUSY flavor problem if the origins of visible flavor structure itself lie at scales *above* v_{mess}. In this paper we are not committing to a theory of how such flavor structure arises in the visible sector, but only mechanisms to avoid the SUSY flavor problem. Therefore the most conservative position is to take as low a v_{mess} as is consistent with other model building constraints, so as to allow for the most "room" in the UV for flavor structure to arise through unspecified dynamics.

3.3 Other dangerous channels

Even though the discrete hidden symmetry is only broken in the IR, the IR-sensitivity of Poppitz-Trivedi mediation (logarithmic in the messenger mass) means the log factor represents an order one breaking of the symmetry. This is an important point which should be compared with another a priori dangerous SUSY-breaking mediation channel. In addition to Planck-suppressed visible-hidden and messenger-hidden Kahler terms, one can also have visible-messenger couplings such as

$$\Delta K = \frac{Q_i^\dagger Q_j (\chi^\dagger \chi + \psi^\dagger \psi)}{M_{Pl}^2}. \tag{3.10}$$

Note that the messenger fields must form an invariant combination under (i) since the visible fields are. Starting with such an interaction and integrating out a messenger loop, threatens to

mediate the SUSY breaking in Eq. (3.8) to the visible sector. The diagrams are UV quadratically divergent and yield visible *flavor-violating* soft mass-squareds $\sim \Lambda_{int}^4 \Lambda_{UV}^2/(16\pi^2 M_{Pl}^4)$. For $M_{Pl} < \Lambda_{UV} < 4\pi M_{Pl}$ this would disasterously dominate over AMSB and GMSB contributions. However, the sum of all such divergent diagrams cancel because of the discrete symmetry (i), since visible fields are even, while the unsequestered messenger mass-squareds are purely odd (originating from the spontaneous breaking of the discrete symmetry by the hidden sector VEVs).

Of course, since the messengers also have (i)-violating supersymmetric masses, there will be less divergent contributions to visible mass-squareds which do not cancel. But they will be suppressed by the messenger masses precisely because highly UV divergent diagrams are IR insensitive, unlike the Poppitz-Trivedi contributions. That is, they will yield visible soft mass-squareds of order $\sim \Lambda_{int}^4 v_{mess}^2/(16\pi^2 M_{Pl}^4)$. For the resulting (generally flavor-violating) SUSY breaking to be as suppressed as the conformally sequestered contributions in Ref. [19], we must have

$$\frac{v_{mess}}{M_{Pl}} < 8 \times 10^{-3}. \tag{3.11}$$

There are also a priori dangerous contributions from Kahler couplings of all three sectors, such as

$$\Delta K \sim \frac{Q_i^\dagger Q_j (|A|^2 - |B|^2)(|T_{1a}|^2 - |T_{2a}|^2 + |T_{3a}|^2 - |T_{4a}|^2)}{M_{Pl}^4}. \tag{3.12}$$

(The lower dimension operator with single powers of A and B could be forbidden by symmetries $A, B, \chi, \overline{\psi} \to -A, B, \chi, \overline{\psi}$ respectively, with other fields inert.) Again, the hidden flavor non-singlet current is not renormalized so that in the IR this will match onto an operator

$$\Delta K \sim \frac{Q_i^\dagger Q_j v_{mess}^2 X^\dagger X}{M_{Pl}^4}. \tag{3.13}$$

Again, this is sufficiently suppressed if Eq. (3.11) is satisified.

4 Realistic Example

Ref. [12] engineered effective field theories in flat 5D compactification with the same interplay of AMSB and Poppitz-Trivedi gauge mediation which we have realized here through (non-perturbative) 4D gauge dynamics. It is important to stress that the mechanisms are not in any sense AdS/CFT dual to each other, indeed the proposal of Ref. [12] does not extend to (truncated) 5D AdS spacetimes. The proposal of the present paper is an independent mechanism demonstrating how AMSB and gauge mediation can naturally combine.

However, the phenomenological results of Ref. [12] can be carried over to the present paper. Ref. [12] chose the visible sector to consist of just the NMSSM, with the hope that the singlet would develop a VEV to act as the μ-term. If AMSB by itself dominated visible SUSY breaking, while all soft masses would be of the same parametric size $\sim$ loop-factor $\times \Lambda_{int}^2/M_{Pl}$, the signs do not work out. Sleptons obtain tachyonic mass-squareds, while the singlet does not, and hence does not develop a weak scale VEV. Running from the Planck scale cannot help because of the characteristic UV insensitivity of AMSB. Fortunately, the gauge-mediated contributions can correct the sign of the slepton mass-squareds and are sensitive to running from the UV. The singlet can then be driven to condense at the weak scale and solve the μ-problem.

Ref. [12] found fully realistic examples in parameter space by treating the gauge-mediated soft Higgs mass-squareds as independent from the gauge-mediated soft slepton mass-squareds. This is only possible at or above the GUT scale where the Higgs and slepton gauge interactions can be distinct, which requires that the messenger scale $v_{mess} \gtrsim M_{GUT}$. This realistic example carries over to the present paper. However, note that the requirement $v_{mess} \gtrsim M_{GUT}$ is only just consistent with the requirement of Eq. (3.11).

5 Extensions

Let us consider why we might wish to pursue extensions of our results so far. While a visible NMSSM sector allows one to find realistic and natural SUSY and electroweak symmetry breaking by combining AMSB and gauge-mediation, as pointed out in Ref. [12] such examples are difficult to find in the visible parameter space. Also, it is rather restrictive that v_{mess} is constrained from above and below to be about the GUT scale in order for the Higgs and slepton gauge-mediated soft mass-squareds to be independent. This restricts visible sector flavor structure to originate in the narrow range of energies between the GUT and Planck scales. Finally, given the mechanism demonstrated here in conformal sequestering that gauge-mediation and AMSB can fruitfully compete, it is worthwhile understanding more general forms of the gauge-mediated soft terms before committing to a particular visible sector model.

It is a straightforward to extend our visible sector from the NMSSM to include a $U(1)'$ gauge group, where the charge is given by a linear combination of hypercharge and $B - L$ which commutes with $SU(5)$, which does not interfere with standard model gauge unification. To cancel anomalies one includes the right-handed neutrino supermultiplet, and to eliminate this gauge group by observable energies one includes fields that can Higgs the $U(1)'$. We can then add extra messengers which are standard model singlets but are vectorially charged under $U(1)'$, with their own Planck-suppressed couplings to the hidden sector. This will result in an independent Poppitz-Trivedi gauge mediation contribution to the visible sector, but now

proportional to powers of $U(1)'$ charges, rather than standard model charges. (See Refs. [27] [28] [29] for earlier proposals of $B-L$ mediation.) In this way the gauge mediated contributions to the Higgs mass-squareds are independent of the gauge mediated contributions to the squarks and sleptons even when $v_{mess} \ll M_{GUT}$. The only requirement on the $U(1)'$ breaking scale is that it be below v_{mess}. Still, this can be at such a high scale that none of the extra $U(1)'$-related structure need survive to the weak scale, where we can have just the NMSSM. The lower v_{mess}, the larger the range of energies at which visible sector flavor structure can emerge.

The basic mechanisms decribed here may be compatible with a variety of visible sectors. Which of these most naturally accomodates the data with minimal fine-tuning at the weak scale remains an important and ongoing investigation. In future work the possibility of naturally combining AMSB with "F-type" GMSB will be explored.

Acknowledgements

I am very grateful to Kaustubh Agashe, David E. Kaplan, Markus Luty and Tsutomu Yanagida for helpful discussions. This work was supported by NSF Grant PHY-0099468.

References

[1] M. Dine, arXiv:hep-ph/9612389.

[2] L. Alvarez-Gaume, M. Claudson and M. B. Wise, Nucl. Phys. B **207**, 96 (1982); M. Dine and W. Fischler, Nucl. Phys. B **204**, 346 (1982).

[3] M. Dine and A. E. Nelson, Phys. Rev. D **48**, 1277 (1993) [arXiv:hep-ph/9303230]; M. Dine, A. E. Nelson and Y. Shirman, Phys. Rev. D **51**, 1362 (1995) [arXiv:hep-ph/9408384]; M. Dine, A. E. Nelson, Y. Nir and Y. Shirman, Phys. Rev. D **53**, 2658 (1996) [arXiv:hep-ph/9507378].

[4] G. F. Giudice and R. Rattazzi, Phys. Rept. **322**, 419 (1999) [arXiv:hep-ph/9801271].

[5] D.E. Kaplan, G.D. Kribs, M. Schmaltz, *Phys. Rev.* **D62** (2000) 035010, hep-ph/9911293; Z. Chacko, M.A. Luty, A.E. Nelson, E. Pontón, *JHEP* **0001** (2000) 003, hep-ph/9911323.

[6] L. Randall, R. Sundrum, *Nucl. Phys.* **B557** (1999) 79, hep-th/9810155.

[7] G.F. Giudice, M.A. Luty, H. Murayama, R. Rattazzi, *JHEP* **9812** (1998) 027, hep-ph/9810442.

[8] A. Pomarol and R. Rattazzi, JHEP **9905**, 013 (1999) [arXiv:hep-ph/9903448].

[9] D. E. Kaplan and G. D. Kribs, JHEP **0009**, 048 (2000) [arXiv:hep-ph/0009195].

[10] I. Jack and D. R. T. Jones, Phys. Lett. B **482**, 167 (2000) [arXiv:hep-ph/0003081].

[11] N. Arkani-Hamed, D. E. Kaplan, H. Murayama and Y. Nomura, JHEP **0102**, 041 (2001) [arXiv:hep-ph/0012103].

[12] Z. Chacko and M. A. Luty, JHEP **0205**, 047 (2002) [arXiv:hep-ph/0112172].

[13] A. E. Nelson and N. J. Weiner, Phys. Rev. Lett. **88**, 231802 (2002) [arXiv:hep-ph/0112210].

[14] E. Poppitz and S. P. Trivedi, Phys. Lett. B **401**, 38 (1997) [arXiv:hep-ph/9703246].

[15] M.A. Luty, R. Sundrum, *Phys. Rev.* **D62** (2000) 035008, hep-th/9910202.

[16] A. Anisimov, M. Dine, M. Graesser and S. Thomas, Phys. Rev. D **65**, 105011 (2002) [arXiv:hep-th/0111235]; A. Anisimov, M. Dine, M. Graesser and S. Thomas, JHEP **0203**, 036 (2002) [arXiv:hep-th/0201256].

[17] A.E. Nelson, M.J. Strassler, `hep-ph/0104051`; *JHEP* **0009** (2000) 030, `hep-ph/0006251`; A. E. Nelson and M. J. Strassler, JHEP **0207**, 021 (2002) [arXiv:hep-ph/0104051].

[18] M. A. Luty and R. Sundrum, Phys. Rev. D **65**, 066004 (2002) [arXiv:hep-th/0105137].

[19] M. Luty and R. Sundrum, Phys. Rev. D **67**, 045007 (2003) [arXiv:hep-th/0111231].

[20] M. A. Luty and R. Sundrum, Phys. Rev. D **64**, 065012 (2001) [arXiv:hep-th/0012158].

[21] J. Maldacena, *Adv. Theor. Math. Phys.* **2** (1998) 231, `hep-th/9711200`; S.S. Gubser, I.R. Klebanov, A.M. Polyakov, *Phys. Lett.* **B428** (1998) 105, `hep-th/9802109`; E. Witten, *Adv. Theor. Math. Phys.* **2** (1998) 253, `hep-th/9802150`.

[22] L. Randall, R. Sundrum, *Phys. Rev. Lett.* **83** (1999) 3370, `hep-ph/9905221`; L. Randall, R. Sundrum, *Phys. Rev. Lett.* **83** (1999) 4690, `hep-th/9906064`.

[23] H. Verlinde, hep-th/9906182, Nucl. Phys. B 580, 264 (2000); J. Maldacena, unpublished remarks; E. Witten, ITP Santa Barbara conference 'New Dimensions in Field Theory and String Theory', `http://www.itp.ucsb.edu/online/susy_c99/discussion`; S. Gubser, hep-th/9912001, Phys. Rev. D 63, 084017 (2001); E. Verlinde and H. Verlinde, hep-th/9912018, JHEP 0005, 034 (2000); N. Arkani-Hamed, M. Porrati and L. Randall, hep-th/0012148, JHEP 0108, 017 (2001); R. Rattazzi and A. Zaffaroni, hep-th/0012248, JHEP 0104, 021 (2001); M. Perez-Victoria, hep-th/0105048, JHEP 0105, 064 (2001).

[24] M. Dine, P. J. Fox, E. Gorbatov, Y. Shadmi, Y. Shirman and S. Thomas, hep-ph/0405159.

[25] R. Harnik, H. Murayama and A. Pierce, JHEP **0208**, 034 (2002) [arXiv:hep-ph/0204122].

[26] J. Polonyi, Budapest preprint no.KFKI-1977-93 (1977).

[27] I. Affleck, M. Dine and N. Seiberg, "Dynamical Supersymmetry Breaking In Four-Dimensions And Its Phenomenological Nucl. Phys. B **256**, 557 (1985).

[28] R. N. Mohapatra and S. Nandi, Phys. Rev. Lett. **79**, 181 (1997) [arXiv:hep-ph/9702291].

[29] B. A. Dobrescu, Phys. Lett. B **403**, 285 (1997) [arXiv:hep-ph/9703390].